

# Higgs Chaotic Inflation and the Primordial B-mode Polarization Discovered by BICEP2

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## Abstract

We show that the standard model Higgs field can realize the quadratic chaotic inflation, if the kinetic term is significantly modified at large field values. This is a simple realization of the so-called running kinetic inflation. The point is that the Higgs field respects an approximate shift symmetry at high energy scale. The tensor-to-scalar ratio is predicted to be  $r \simeq 0.13-0.16$ , which nicely explains the primordial B-mode polarization,  $r = 0.20^{+0.07}_{-0.05}$ , recently discovered by the BICEP2 experiment. In particular, allowing small modulations induced by the shift symmetry breaking, the negative running spectral index can also be induced. The reheating temperature is expected to be so high that successful thermal leptogenesis is possible. The suppressed quartic coupling of the Higgs field at high energy scales may be related to the Higgs chaotic inflation.

Our Universe experienced an accelerated expansion at a very early stage of the evolution, i.e., inflation [1, 2]. Among various inflation models proposed so far, the so-called chaotic inflation [3] is particularly interesting as it predicts large values of the tensor-to-scalar ratio  $r$ . The tensor mode density perturbations generate the primordial B-mode polarization of the cosmic microwave background (CMB), which, if observed, would determine the inflation scale and pin down the underlying model of inflation.

Recently the BICEP2 experiment announced that they discovered the primordial B-mode polarization of CMB. In terms of the tensor-to-scalar ratio, the allowed range is given by [4]

$$r = 0.20^{+0.07}_{-0.05} \quad (68\% \text{CL}). \quad (1)$$

After subtracting the best available estimate for foreground dust, the allowed range is modified to  $r = 0.16^{+0.06}_{-0.05}$ . The discovery of  $r$  in this range is of significant importance for cosmology as well as particle physics, as it implies that we obtain the invaluable information on the Universe at the GUT scale.

There are various large-field inflation models which predict  $r$  in the above range<sup>1</sup>, and by far the simplest one is the quadratic chaotic inflation [3]:

$$\mathcal{L} = \frac{1}{2}(\partial\phi)^2 - \frac{1}{2}m^2\phi^2, \quad (2)$$

where  $\phi$  is the inflaton and  $m$  is the inflaton mass. The Planck normalization on the primordial density perturbations [17] fixes the inflaton mass as

$$m \simeq 1.5 \times 10^{13} \text{ GeV}. \quad (3)$$

Since the primordial B-mode polarization was discovered, the next question will be the identity of the inflaton.

In fact, there is a unique scalar field in the standard model (SM), i.e., the Higgs field, which was discovered at LHC in 2012 [18, 19].<sup>2</sup> In order for the Higgs field to realize the

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<sup>1</sup> For various large-field inflation models and their concrete realization in supergravity and superstring theory, see e.g. [5, 6, 7, 8, 9, 10, 11, 12, 13, 14, 15, 16].

<sup>2</sup> The connection between the SM Higgs field and inflation has been extensively discussed especially in a context of the non-minimal coupling to gravity [20, 21, 22, 23, 24, 25]. The Starobinsky-type inflation, however, leads to much smaller values of  $r$ .

quadratic chaotic inflation with correct density perturbations, however, the Higgs field must have a mass of order  $10^{13}$  GeV, many orders of magnitude larger than the observed mass,  $m_h \approx 126$  GeV. Moreover, the Higgs potential is dominated by the quartic term at large field values, not the quadratic one. The apparent discrepancy can be reconciled if either the kinetic term or the potential term is modified at large field values. Actually, the present authors proposed in Ref. [12] a Higgs chaotic inflation model in which the SM Higgs field realizes the quadratic chaotic inflation model, based on the so-called running kinetic inflation [10, 11]. In this letter we revisit the SM Higgs chaotic inflation model, in light of the recent discoveries of the SM Higgs boson as well as the primordial CMB B-mode polarization.

The basic idea of the running kinetic inflation is very simple. Let us consider a scalar field with the following Lagrangian,

$$\mathcal{L} = \frac{1}{2} (1 + \xi \phi^2) (\partial\phi)^2 - V(\phi), \quad (4)$$

where  $\xi$  is a positive numerical coefficient much larger than unity, and  $V(\phi)$  is the inflaton potential. Here and in what follows we adopt the Planck units in which the reduced Planck scale  $M_P \simeq 2.4 \times 10^{18}$  GeV is set to be unity. Due to the dependence of the kinetic term on  $\phi^2$ , the canonically normalized field at  $\phi \gtrsim 1/\sqrt{\xi}$  is given by  $\hat{\phi} \sim \sqrt{\xi} \phi^2$ . As the kinetic term grows, the potential in terms of the canonically normalized field becomes flatter. For instance, the quartic potential,  $V(\phi) \sim \phi^4$ , becomes the quadratic one,  $V(\hat{\phi}) \sim \hat{\phi}^2/\xi$ , at large field values. This is the essence of the running kinetic inflation. The running kinetic inflation can be easily implemented in supergravity and the cosmological implications were studied in Refs. [10, 11].

The above argument can be straightforwardly applied to the SM Higgs field, and the SM Higgs can drive quadratic chaotic inflation [12]. In order to build sensible inflation models, we need to have a good control of the scalar potential over large field values. Also, it is desirable to understand the large value of  $\xi \gg 1$  in terms of symmetry. To this end, we impose an approximate shift symmetry on absolute square of the SM Higgs field  $H$  [10, 11, 12]:

$$|H|^2 \rightarrow |H|^2 + C, \quad (5)$$

where  $C$  is a real parameter and the  $SU(2)_L$  indices are suppressed. We assume that the shift symmetry exhibits itself at high energy scales, whereas it is explicitly broken and therefore becomes much less prominent at low energy scales. Then we can write down the Lagrangian of the Higgs field at high energy scales as follows:<sup>3</sup>

$$\mathcal{L} = \frac{1}{2} (\partial_\mu |H|^2)^2 + \epsilon |D_\mu H|^2 - \lambda \left( |H|^2 - \frac{v^2}{2} \right)^2 + \cdots, \quad (6)$$

where  $\epsilon$  and  $\lambda$  are coupling constants, and  $D_\mu$  denotes the covariant derivative. The first term in (6) respects the shift symmetry, which is explicitly broken by the second and third terms, and so, we expect  $\epsilon, \lambda \ll 1$ . Note that the second term provides the usual kinetic term for the SM Higgs in the low energy, whereas the first term provides the kinetic term for  $|H|^2$  at large field values. In the unitary gauge, the relevant interactions are<sup>4</sup>

$$\mathcal{L} = \frac{1}{2} (\epsilon + h^2) (\partial h)^2 - \frac{\lambda}{4} (h^2 - v^2)^2, \quad (7)$$

where  $h$  denotes the physical Higgs boson, and we have omitted the gauge and Yukawa interactions which are irrelevant during inflation. Thus, the largeness of  $\xi$  in Eq. (4) is due to the smallness of  $\epsilon$ , i.e., the fact that the usual kinetic term breaks the shift symmetry (5).

For large field value  $h \gg \sqrt{\epsilon}$ , we can rewrite the Lagrangian in terms of the canonically normalized field  $\hat{h} \equiv h^2/2$  as

$$\mathcal{L} \simeq \frac{1}{2} (\partial \hat{h})^2 - \lambda \hat{h}^2. \quad (8)$$

Thus we obtain the quadratic potential in terms of  $\hat{h}$ . The Planck normalization on the density perturbation (3) fixes  $\lambda = m^2/2 \simeq 2 \times 10^{-11}$ . Thus, the chaotic inflation with quadratic potential can be realized by the SM Higgs field with the running kinetic term. Note that all the interaction of the (canonically normalized) Higgs field are suppressed and the system approaches the free field theory as  $h$  increases.

The BICEP2 result (1) has an apparent tension with the Planck data, which can be relaxed by including a large negative running spectral index [4]. A recent work [27] have

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<sup>3</sup> In general, we can add an arbitrary Lorentz-invariant function of  $\partial|H|^2$ , which preserves the shift symmetry, as well as other interactions of the Higgs fields, which break the shift symmetry.

<sup>4</sup> The Higgs chaotic inflation with the same Lagrangian was also studied in Ref. [26].

performed a joint analysis of the *Planck* and BICEP2 datasets, giving a constraint,

$$\frac{dn_s}{d\ln k} = -0.024 \pm 0.010 \quad (68\% \text{CL}), \quad (9)$$

which is similar to the combined analysis of *Planck*+WP+highL data [28]. The analyses assume a scale-independent running, and it is known that such a large (constant) negative running would quickly terminate inflation within the e-folding number 30 or so [29]. This conclusion, however, can be avoided by allowing a scale-dependence of the running, while it remains more or less constant over the CMB scales. The simplest way to accomplish this is to add small modulations to the inflaton potential [30]. The point is that the third derivative of the inflaton potential, which contributes to the running, can be (locally) dominated by the modulations, while their contributions to the potential and its first derivative are negligibly small. Therefore, the running can be enhanced locally without modifying the overall inflaton dynamics. In our case, we can add small modulations to the Higgs potential by introducing another shift symmetry breaking term<sup>5</sup>

$$\mathcal{L} \supset \Lambda^4 \cos(|H|^2/f + \theta). \quad (10)$$

The precise value of the running depends on the phase of modulations, but we can evaluate the typical value of the running as

$$\left| \frac{dn_s}{d\ln k} \right| \simeq \left| -2 \frac{V'V'''}{V^2} \right| \sim \frac{\Lambda^4}{N^{3/2} m^2 f^3}, \quad (11)$$

where  $N$  is the e-folding number. The running spectral index (9) can be realized for e.g.  $N = 60$ ,  $f \sim 0.1$  and  $\Lambda^4 \sim m^2 f^2$  [31].

For small field value  $h \ll \sqrt{\epsilon}$ , the Lagrangian is reduced to the usual one for the SM Higgs field,

$$\mathcal{L} \simeq \frac{1}{2}(\partial\tilde{h})^2 - \frac{\tilde{\lambda}}{4} \left( \tilde{h}^2 - \tilde{v}^2 \right)^2, \quad (12)$$

where we have defined  $\tilde{h} \equiv \sqrt{\epsilon}h$ ,  $\tilde{\lambda} \equiv \lambda/\epsilon^2$  and  $\tilde{v} \equiv \sqrt{\epsilon}v$ . In order to explain the correct electroweak scale and the 126 GeV Higgs boson mass, we must have  $\tilde{v} = 246$  GeV and  $\tilde{\lambda} \simeq 0.13$ . The SM Yukawa interactions are also obtained in the low energy if we add

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<sup>5</sup> This additional interactions will contribute to the Higgs potential in the low energy, but its effect can be absorbed by shifting the values of  $\lambda$  and  $v$ .

the Yukawa interactions with suppressed couplings in Eq. (6) [12]. That is to say, the low-energy effective theory coincides with the SM at  $\tilde{h} \lesssim \epsilon$ , while the theory approaches the free theory for  $\tilde{h}$  at  $\tilde{h} \gtrsim \epsilon$ . The transition from one phase to the other takes place at the intermediate scale  $\sim 10^{13}$  GeV or above, whose precise value depends on the running of the quartic coupling  $\tilde{\lambda}$ , as shown below.

The quartic coupling  $\tilde{\lambda}$  evolves through the renormalization group equations (RGE), and it is known that the quartic coupling and its beta function become tiny at high energy scale [32, 33, 34]. Allowing the RGE running of  $\tilde{\lambda}$  in the SM regime, we can estimate  $\epsilon$  as

$$\epsilon \simeq 1 \times 10^{-5} \sqrt{\frac{\tilde{\lambda}_{\text{IR}}}{\tilde{\lambda}_{\text{UV}}}} \left( \frac{\lambda}{2 \times 10^{-11}} \right)^{\frac{1}{2}}, \quad (13)$$

where  $\tilde{\lambda}_{\text{UV}}$  and  $\tilde{\lambda}_{\text{IR}}$ , respectively, denote the SM Higgs quartic coupling evaluated at the UV and IR energy scales in the scheme of  $h < \sqrt{\epsilon}$ . That is to say,  $\tilde{\lambda}_{\text{UV}}$  is defined by  $\lambda/\epsilon^2$ , whereas  $\tilde{\lambda}_{\text{IR}}$  is fixed by the Higgs boson mass to be  $\tilde{\lambda}_{\text{IR}} \simeq 0.13$ . Therefore, the smallness of the Higgs quartic coupling at high energy scales,  $\tilde{\lambda}_{\text{UV}} < \tilde{\lambda}_{\text{IR}}$ , is related to the size of  $\epsilon$ , which parametrizes the explicit breaking of the shift symmetry. As  $\tilde{\lambda}_{\text{UV}} < \tilde{\lambda}_{\text{IR}}$  is suggested by the top mass and the Higgs boson mass, the transition from the SM to the free theory occurs above the intermediate scale, i.e.,  $\tilde{h} \sim \epsilon \gtrsim 10^{13}$  GeV.

After inflation ends, the Higgs field begins coherent oscillations. As we have seen, while the Yukawa and gauge interactions of the Higgs field are suppressed at  $\tilde{h} \gtrsim \epsilon$ , the SM interactions are reproduced at  $\tilde{h} \lesssim \epsilon$ . Thus the particle production during the coherent oscillations is considered to be so efficient that the Higgs bosons are thermalized soon. The reheating temperature therefore tends to be very high, and it could be as high as  $\sim 10^{13-14}$  GeV [35]. Thermal leptogenesis works for such a high reheating temperature [36]. With the large tensor-to-scalar ratio  $r$  given by Eq. (1), the reheating temperature may be probed by the future gravitational wave experiments [37, 38].

It is possible to extend our Higgs chaotic inflation model to the linear or fractional power potential by imposing a shift symmetry on a certain combination of the Higgs field. The Higgs chaotic inflation can also be implemented in supergravity [10, 11, 12]. In particular, we can identify the D-flat direction  $H_u H_d$  as the inflaton, once we impose a shift symmetry as  $H_u H_d \rightarrow H_u H_d + C$  [11, 12]. In the simplest case, the Kähler potential

includes a term like,  $K \supset (H_u H_d - (H_u H_d)^\dagger)^2$ , in addition to the usual kinetic terms which break the shifts symmetry, and the scalar potential of  $H_u H_d$  is generated by the superpotential  $W = \lambda S H_u H_d$ , where  $S$  is the gauge singlet field.<sup>6</sup> The inflaton dynamics is the same as the model of (6). By further imposing  $Z_n$  symmetry on  $H_u H_d$  with  $n$  being an integer, Higgs chaotic inflation with fractional power potential can also be realized [12].

There is no candidate for dark matter in the SM. One of the plausible dark matter candidates in the extension of SM is the QCD axion [40, 41]. In the present setup, both the inflation scale and the reheating temperature are so high that the Peccei-Quinn (PQ) symmetry is expected to be restored during or after inflation. As the PQ symmetry gets spontaneously broken after inflation, the axionic strings as well as domain walls appear. To avoid the domain wall problem, the domain wall number must be equal to one. In this case, the axions radiated from the collapse of axionic domain walls gives the dominant contribution to the relic axion abundance, which requires  $f_a \lesssim 3 \times 10^{10} \text{ GeV}$  where  $f_a$  denotes the PQ symmetry breaking scale [42]. Another candidate is the sterile neutrino [43, 44, 45], if it is sufficiently light and long-lived. The light mass can be realized by e.g. the split seesaw mechanism [46], and both the light mass and the longevity can also be explained by the flavor symmetry [47]. Interestingly, the reheating temperature of order  $10^{13-14} \text{ GeV}$  is sufficient for producing the right amount of the sterile neutrinos through the  $B - L$  gauge boson exchange, to account for the observed dark matter density [46]. At the same time, thermal leptogenesis with two heavy right-handed neutrinos is also possible with this temperature [49, 50]. The sterile neutrino dark matter with mass of about 7 keV [48, 51] can explain the recent hint for the 3.5 keV X-ray line [52, 53].

In this letter we have revisited the SM Higgs chaotic inflation in light of the recent discoveries of the SM Higgs at LHC and the primordial B-mode polarization by the BICEP2 experiment, and shown that the quadratic chaotic inflation can be realized by the SM Higgs field, based on the running kinetic inflation. One of the essential ingredients is the shift symmetry of the Higgs field (5). If small modulations to the inflaton potential are induced by the shift symmetry breaking, a sizable running spectral index can also be generated without significant effect on the overall inflaton dynamics. If the primordial B-mode polarization is more precisely measured by the Planck and other ground-based

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<sup>6</sup>See Ref. [39] for a shift symmetry on  $H_u$  and  $H_d$ , instead of on  $H_u H_d$ .

observations, it will pin down the underlying inflation model. Then the next question will be what the inflaton is, which will be important not only for the UV theory but also for considering thermal history of the Universe as the baryon asymmetry and dark matter abundance crucially depend on the reheating temperature in many scenarios. The SM Higgs boson has two important advantages with respect to other candidates for the inflaton. First, it was already discovered and we know that it exists. Second, the reheating takes place through the SM interactions, and there is no need to introduce additional interactions between the inflaton and the SM sector. As to the second point, one can also straightforwardly apply the same argument to the  $B - L$  Higgs field. Then, the non-thermal leptogenesis will be possible. The PQ field can also be the inflaton: it couples to the PQ quarks as well as gluons and hence the reheating takes place successfully. We also note that the running kinetic inflation can be applied to a gauge singlet inflaton.

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